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Observable seesaw and its collider signatures

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ABSTRACT

We discuss the scenario with TeV-scale right-handed neutrinos, which are accessible at future colliders, while holding down tiny seesaw-induced masses and sizable couplings to the standard-model particles. The signal with tri-lepton final states and large missing transverse energy is appropriate for studying collider signatures of the scenario with extra spatial dimensions. We show that the LHC experiment generally has a potential to discover the signs of extra dimensions and the origin of small neutrino masses.

1. Introduction

The recent neutrino oscillation experiments have been revealing the detailed structure of leptonic flavors [1,2]. The neutrino property, in particular the tiny mass scale is one of the most important experimental clues to find the new physics beyond the Standard Model (SM). The seesaw mechanism naturally leads to small neutrino masses by the integration of new heavy particles which interact with the ordinary neutrinos. The introduction of heavy right-handed neutrinos [3] implies the intermediate mass scale of such states to have light Majorana masses of order eV, and these heavy states are almost decoupled in the low-energy effective theory. Alternatively, TeV-scale right-handed neutrinos could also be possible, which in turn means tiny orders of couplings to the SM sector and their signs cannot be observed in near future TeV-scale particle experiments such as the Large Hadron Collider (LHC).

The SM neutrinos have tiny masses due to a slight violation of the lepton number. This fact implies that the events with same-sign di-lepton final states [4] are too rare to be observed. In this Letter, we focus on the lepton number preserving processes, in particular, the tri-lepton signals with large missing transverse energy, $pp \to \ell^\pm\ell^\mp\ell^\pm\nu(\bar{\nu})$. These processes would be rather effectively detected at the LHC because only small fraction of SM processes contributes to the background against the signals.

As a simple example of observable seesaw theory, we consider a five-dimensional extension of the SM with right-handed neutrinos, where all SM fields are confined in the four-dimensional world, while right-handed neutrinos propagate in the whole extra-dimensional space [5–7]. We will discuss an explicit framework which provides the situation that TeV-scale right-handed neutrinos generate tiny scale of seesaw-induced neutrino masses and simultaneously have sizable interactions to the SM leptons and gauge bosons. The scenario does not rely on any particular generation structure of mass matrices and is available for one-generation case. For such TeV-scale particles with large couplings to the SM sector, the LHC experiment generally has the potential to find the signals of extra dimensions and the origin of small neutrino masses.

2. Observable seesaw

Let us consider a five-dimensional theory where the extra space is compactified on the S^1/Z_2 orbifold with the radius R. The SM fields are confined on the four-dimensional boundary at $x^5 = 0$. Besides the gravity, only SM gauge singlets can propagate in the bulk not to violate the charge conservation [5,6]. The gauge-singlet Dirac fermions $\mathcal{N}_i$ $(i = 1, 2, 3)$ are introduced in the bulk which contain the right-handed neutrinos and their chiral partners. The Lagrangian up to the quadratic order of spinor fields is given by

$$\mathcal{L} = i\bar{\mathcal{N}}\not{D}\mathcal{N} - \frac{1}{2}\left[\bar{\mathcal{N}}^c(M_v + M_a\gamma_5)\mathcal{N} + \text{h.c.}\right]. \tag{2.1}$$

The conjugated spinor is defined as $\mathcal{N}^c = \gamma_3\gamma_1\bar{\mathcal{N}}^{\rm t}$ such that it is Lorentz covariant in five dimensions. It is straightforward to write a bulk Dirac mass for $\mathcal{N}_i$ if introducing a Z_2-odd function which originates from some field expectation value. The bulk mass parameters M_v and M_a are Z_2 parity even and could depend on the extra-dimensional coordinate x^5 which comes from the delta-function dependence (resulting in localized mass terms) and/or the

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background geometry such as the warp factor in AdS$_5$. We also introduce the mass terms between bulk and boundary fields:

$$\mathcal{L}_m = -\big(\bar{\mathcal{N}} m L + \bar{\mathcal{N}}^c m^c L\big)\delta\big(x^5\big) + \text{h.c.}, \tag{2.2}$$

where m and m^c denote the mass parameters after the electroweak symmetry breaking. The boundary spinors L_i $(i = 1, 2, 3)$ contain the left-handed neutrinos ν_i, namely, given in the 4-component notation $L_i = \binom{0}{\nu_i}$. The Z_2 parity implies that either component in a Dirac fermion $\mathcal{N}$ vanishes at the boundary ($x^5 = 0$) and therefore either of m and m^c becomes irrelevant.[1] In the following we assign the even Z_2 parity to the upper (right-handed) component of bulk fermions, i.e. $\mathcal{N}(-x^5) = \gamma_5 \mathcal{N}(x^5)$, and will drop the m^c term.

With a set of boundary conditions, the bulk fermions $\mathcal{N}_i$ are expanded by Kaluza–Klein (KK) modes with their kinetic terms being properly normalized

$$\mathcal{N}\big(x, x^5\big) = \begin{pmatrix} \sum_n \chi_R^n(x^5) N_R^n(x) \\ \sum_n \chi_L^n(x^5) N_L^n(x) \end{pmatrix}. \tag{2.3}$$

The wavefunctions $\chi_{R,L}^n$ are generally matrix-valued in the generation space and we have omitted the generation indices for notational simplicity. After integrating over the fifth dimension, we obtain the neutrino mass matrix in four-dimensional effective theory. Neutrinos are composed of the boundary ones and the KK modes $(\nu, \epsilon N_R^{0*}, \epsilon N_R^{1*}, N_L^1, \ldots) \equiv (\nu, N)$. The zero modes of the left-handed components have been extracted according to the boundary condition. The neutrino mass matrix for (ν, N) is given by

$$\left(\begin{array}{c|cccc} & m_0^{\mathrm t} & m_1^{\mathrm t} & 0 & \cdots \\ \hline m_0 & M_{R_{00}}^* & M_{R_{01}}^* & M_{K_{01}} & \cdots \\ m_1 & M_{R_{10}}^* & M_{R_{11}}^* & M_{K_{11}} & \cdots \\ 0 & M_{K_{10}}^{\mathrm t} & M_{K_{11}}^{\mathrm t} & M_{L_{11}} & \cdots \\ \vdots & \vdots & \vdots & \vdots & \ddots \end{array}\right) \equiv -\left(\begin{array}{c|c} & M_D^{\mathrm t} \\ \hline M_D & M_N \end{array}\right), \tag{2.4}$$

where the boundary Dirac masses m_n, the KK masses M_K, and the Majorana masses $M_{R,L}$ are

$$\begin{aligned} m_n &= \chi_R^{n\dagger}(0) m, \\ M_{R_{mn}} &= \int_{-\pi R}^{\pi R} dx^5 \, \big(\chi_R^m\big)^{\mathrm t} (M_a + M_v) \chi_R^n, \\ M_{K_{mn}} &= \int_{-\pi R}^{\pi R} dx^5 \, \big(\chi_R^m\big)^{\dagger} \partial_5 \chi_L^n, \\ M_{L_{mn}} &= \int_{-\pi R}^{\pi R} dx^5 \, \big(\chi_L^m\big)^{\mathrm t} (M_a - M_v) \chi_L^n. \end{aligned} \tag{2.5}$$

It is noticed that $M_{K_{mn}}$ becomes proportional to δ_{mn} if $\chi_{R,L}^n$ are the eigenfunctions of the bulk equations of motion, and $M_{R,L_{mn}}$ also becomes proportional to δ_{mn} if the bulk mass parameters M_a, M_v are independent of the coordinate x^5.

We further implement the seesaw operation assuming $\mathcal{O}(m_n) \ll \mathcal{O}(M_K)$ or $\mathcal{O}(M_{L,R})$ and find the induced Majorana mass matrix for three-generations light neutrinos

$$M_\nu = M_D^{\mathrm t} M_N^{-1} M_D. \tag{2.6}$$

It is useful for later discussion of collider phenomenology to write down the electroweak Lagrangian in the basis where all the mass matrices are generation diagonalized. The interactions to the electroweak gauge bosons are given in this mass eigenstate basis (ν_d, N_d) as follows:

$$\begin{aligned} \mathcal{L}_g = {}& \frac{g}{\sqrt{2}} \big[W_\mu^\dagger e^\dagger \sigma^\mu U_{\text{MNS}} (\nu_d + V N_d) + \text{h.c.} \big] \\ & + \frac{g}{2\cos\theta_W} Z_\mu \big(\nu_d^\dagger + N_d^\dagger V^\dagger\big) \sigma^\mu (\nu_d + V N_d), \end{aligned} \tag{2.7}$$

where W_μ and Z_μ are the electroweak gauge bosons and g is the $SU(2)_{\text{weak}}$ gauge coupling constant. The 2-component spinors ν_d are three light neutrinos for which the seesaw-induced mass matrix M_ν is diagonalized

$$M_\nu = U_\nu^* M_\nu^d U_\nu^\dagger, \qquad U_\nu \nu_d = \nu - M_D^\dagger M_N^{-1*} N, \tag{2.8}$$

and N_d denote the infinite number of neutrino KK modes for which the bulk mass matrix M_N is diagonalized in the generation and KK spaces by a unitary matrix U_N:

$$M_N = U_N^* M_N^d U_N^\dagger, \qquad U_N N_d = N + M_N^{-1} M_D \nu. \tag{2.9}$$

The lepton mixing matrix measured in the neutrino oscillation experiments is given by $U_{\text{MNS}} = U_e^\dagger U_\nu$, where U_e is the left-handed rotation matrix for diagonalizing the charged-lepton Dirac masses. It is interesting to find that the model-dependent parts of electroweak gauge vertices are governed by a single matrix V which is defined as

$$V = U_\nu^\dagger M_D^\dagger M_N^{-1*} U_N. \tag{2.10}$$

When one works in the basis where the charged-lepton sector is flavor diagonalized, U_ν is fixed by the neutrino oscillation matrix.

The neutrinos also have the interaction to the electroweak doublet Higgs H in four dimensions. The boundary Dirac mass (2.2) comes from the Yukawa coupling

$$\mathcal{L}_h = -\big(y \bar{\mathcal{N}} L H^\dagger + \text{h.c.}\big)\delta\big(x^5\big). \tag{2.11}$$

The doublet Higgs H has a non-vanishing expectation value v and its fluctuation $h(x)$. After integrating out the fifth dimension and diagonalizing mass matrices, we have

$$\mathcal{L}_h = \frac{-1}{v} \sum_n \big[\big(N_d^{\mathrm t} - \nu_d^{\mathrm t} V^*\big) U_N^{\mathrm t}\big]_{R_n} m_n U_\nu \epsilon (\nu_d + V N_d) h^* + \text{h.c.}, \tag{2.12}$$

where $[\cdots]_{R_n}$ means the nth mode of the right-handed component.

The heavy neutrino interactions to the SM fields are determined by the mixing matrix V both in the gauge and Higgs vertices. The $3 \times \infty$ matrix V is determined by the matrix forms of neutrino masses in the original Lagrangian $\mathcal{L} + \mathcal{L}_b$. The matrix elements in V have the experimental upper bounds from electroweak physics, as will be seen later. Another important constraint on V comes from the low-energy neutrino experiments, namely, the seesaw-induced masses should be of the order of eV scale, which in turn specifies the scale of heavy neutrino masses M_N. This can be seen from the definition of V by rewriting it with the light and heavy neutrino mass eigenvalues

$$V = \big(M_\nu^d\big)^{\frac{1}{2}} P \big(M_N^d\big)^{-\frac{1}{2}}, \tag{2.13}$$

where P is an arbitrary $3 \times \infty$ matrix with $PP^{\mathrm t} = 1$. Therefore one naively expects that, with a fixed order of $M_\nu^d \sim 10^{-1}$ eV and $V \gtrsim 10^{-2}$ for the discovery of experimental signatures of heavy neutrinos, their masses should be very light and satisfy $M_N^d \lesssim$ keV (this does not necessarily mean the seesaw operation

[1] The exception is the generation-dependent Z_2 parity assignment on bulk fermions [8]. We do not consider such a possibility in this Letter.

is not justified as M_ν^d is fixed). The previous collider studies on TeV-scale right-handed neutrinos [9] did not impose the seesaw relation (2.13) and have to rely on some assumptions for suppressing the necessarily induced masses M_ν. For example, the neutrino mass matrix has some singular generation structure, otherwise it leads to the decoupling of seesaw neutrinos from collider physics.

A possible scenario for observable heavy neutrinos is to take a specific value of bulk Majorana masses. Here we assume that bulk Dirac masses vanish but it is easy to include them by attaching wavefunction factors in the following formulas. The equations of motion without bulk Majorana masses are solved by simple oscillators and the mass matrices in four-dimensional effective theory are found

$$m_n = \frac{m}{\sqrt{2^{\delta_{n0}}\pi R}}, \qquad M_{R_{mn}} = \delta_{mn}(M_a + M_\nu),$$
$$M_{K_{mn}} = \frac{n}{R}\delta_{mn}, \qquad M_{L_{mn}} = \delta_{mn}(M_a - M_\nu). \tag{2.14}$$

From these, we find the seesaw-induced mass matrix and the mixing with heavy modes:

$$M_\nu = \frac{1}{2\pi R} m^{\mathrm{t}} \frac{\pi R X}{\tan(\pi R X)} \frac{1}{(M_a + M_\nu)^*} m, \tag{2.15}$$

$$\nu = U_\nu \nu_d + \frac{1}{\sqrt{2\pi R}} m^\dagger \left[\frac{1}{M_a + M_\nu} \epsilon N_R^{0*} + \sum_{n=1} \frac{\sqrt{2}}{X^{2*} - (n/R)^2} \left[(M_a - M_\nu)^* \epsilon N_R^{n*} + \frac{n}{R} N_L^n \right] \right], \tag{2.16}$$

where $X^2 = (M_a + M_\nu)^*(M_a - M_\nu)$. The effect of infinitely many numbers of KK neutrinos appears as the factor $\tan(\pi R X)$. An interesting case is that (the eigenvalue(s) of) X takes a specific value $X \simeq \alpha/R$ where α contains half integers [5]: the seesaw-induced mass matrix M_ν is suppressed by the tangent factor (not only by a large Majorana mass scale), on the other hand, the heavy mode interaction V is un-suppressed. This fact realizes the situation that right-handed neutrinos in the seesaw mechanism are observable at sizable rates in future collider experiments.

3. Collider signatures

One of the most exciting signals of higher-dimensional theory at collider experiments is the production of KK excited states. The signals could be observed at the LHC if new physics, which is responsible for the generation of neutrino masses, lies around the TeV scale, and large Yukawa couplings are allowed that lead to a sizable order of mixing between the left- and right-handed neutrinos. An immediate question is what processes we should pay attention to find out the signals. One important possibility is the like-sign di-leptons signal, $pp \to \ell^+ N \to \ell^\pm \ell^\pm W^\mp \to \ell^\pm \ell^\pm jj$, because the SM background against the signal is enough small. Unfortunately, this process violates the lepton number which should be proportional to tiny Majorana neutrino masses, and is therefore difficult to be observed at the LHC. In this Letter we thus focus on lepton number preserving processes. While there are various types of such processes related to heavy neutrino productions, most of these would not be observable due to huge SM backgrounds. As we will see in the following, an exception suitable for the present purpose is the tri-lepton signal with large missing transverse energy: $pp \to \ell^\pm N \to \ell^\pm \ell^\mp W^\pm \to \ell^\pm \ell^\mp \ell^\pm \nu(\bar\nu)$ and $pp \to \ell^\pm N \to \ell^\pm \nu(\bar\nu) Z \to \ell^\pm \nu(\bar\nu) \ell^\pm \ell^\mp$ (Fig. 1). They are possibly captured at the LHC since only small fractions of SM processes contribute to the background against the signal.

To investigate the signal quantitatively, we consider the five-dimensional seesaw theory as a simple example for providing realistic seesaw neutrino masses and observable collider signatures. The right-handed Majorana masses are $M_a = M$ and $M_\nu = 0$ and diagonalized in the generation space. In this Letter it is assumed that these masses are also generation independent. As mentioned before, the effective neutrino Majorana masses become tiny for $M \simeq 1/(2R)$, and thus, the right-handed neutrino masses can be $M \sim 1/R \sim \mathcal{O}(\mathrm{TeV})$, while keeping a non-negligible order of Yukawa couplings and sizable electroweak gauge vertices for the heavy KK neutrinos. We parametrically introduce a small quantity δ as

$$M = \frac{1-\delta}{2R}. \tag{3.1}$$

Summing up the effects of heavy neutrinos,[2] we obtain the seesaw-induced mass $M_\nu = \frac{\delta\pi^2}{8}\frac{m^{\mathrm{t}} m}{M}$. A vanishing value of δ makes the light neutrinos exactly massless, where the complete cancellation occurs in the effects of heavy neutrinos which exhibit the Dirac nature in this case. As we will see, the parameter δ takes a tiny value for giving the right neutrino mass scale. It is noted that the seesaw mechanism generally needs some smallness of (the ratios of) model parameters. That is true for higher-dimensional theory, e.g. a tiny compactification scale in the large extra dimension scenario and also for Majorana neutrinos in warped extra dimension. The present model is an example and there are many other possibilities of the observable seesaw with extra dimensions where model parameters are not finely tuned. We will discuss these issues in a separate paper. The nth excited KK mode spectrum becomes $M_n = (2n-1)/(2R)$.

The electroweak gauge and Higgs vertices are also evaluated from the Lagrangian given in the previous section. For example, the neutrino Yukawa matrix y in the model is expressed as

$$\frac{y}{\sqrt{2\pi R}} = \frac{2}{\pi v}\frac{1}{\sqrt{\delta R}} O^\dagger (M_\nu^d)^{\frac{1}{2}} U_{\mathrm{MNS}}^\dagger, \tag{3.2}$$

where O is the 3×3 orthogonal matrix, which generally comes in reconstructing high-energy quantities from the observable ones [11]. That corresponds to the matrix P in (2.13). The model therefore contains the parameters R, δ, M_ν^d, U_{MNS}, and O. The neutrino mass differences and the generation mixing parameters have been measured and we take their typical experimental values [2]: $\Delta m_{21}^2 = 8 \times 10^{-5}\ \mathrm{eV}^2$, $\Delta m_{32}^2 = 2.5 \times 10^{-3}\ \mathrm{eV}^2$, $\sin\theta_{12} = 0.56$, $\sin\theta_{23} = 0.71$, and $|\sin\theta_{13}| \leqslant 0.22$. In this Letter we consider the neutrino mass spectrum with the normal hierarchy. The other cases of the inverted and degenerate mass patterns can also be analyzed in similar fashion. The Majorana phases in U_{MNS} have no physical relevance in the present work and are set to be zero. The remaining quantities suffer from experimental constraints in low-energy physics. In particular, the dominant constraint is found to come from the experimental search for lepton flavor-changing processes [12,13]. For a real orthogonal matrix O, the limits imposed by lepton flavor conservation are summarized as

$$\frac{2R}{\delta} U_{\mathrm{MNS}} M_\nu U_{\mathrm{MNS}}^\dagger \leqslant \begin{pmatrix} 10^{-2} & 7 \times 10^{-5} & 1.6 \times 10^{-2} \\ 7 \times 10^{-5} & 10^{-2} & 10^{-2} \\ 1.6 \times 10^{-2} & 10^{-2} & 10^{-2} \end{pmatrix}, \tag{3.3}$$

which shows that the most severe limit is given by the 1–2 component, i.e. the $\mu \to e\gamma$ search. We fix $\sin\theta_{13} = 0.07$ as a typical example, and accordingly the Dirac CP phase in U_{MNS} is $\phi_D = \pi$

[2] In theory with more than one extra dimensions, the sums of infinite KK modes generally diverse without regularization [10].

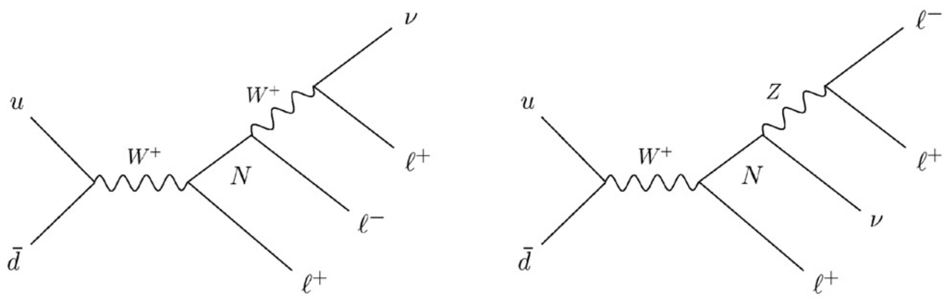

Fig. 1. Lepton number preserving tri-lepton processes at the LHC.

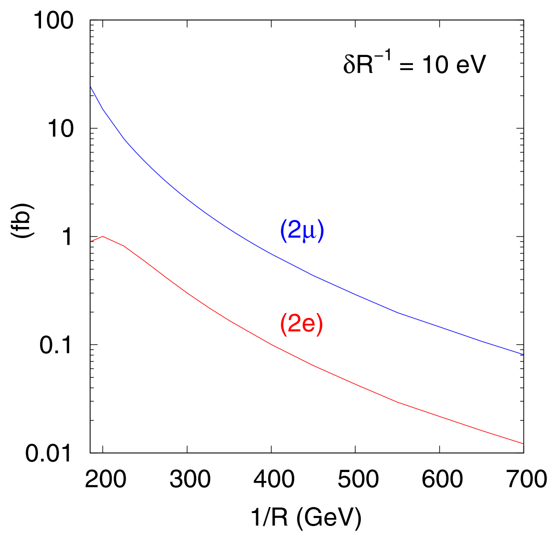

Fig. 2. Total cross sections of tri-lepton signals as the functions of the compactification scale R with a fixed value $\delta/R = 10$ eV.

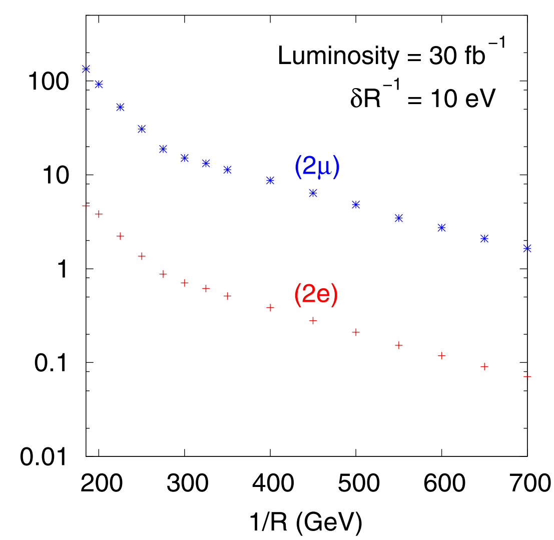

Fig. 3. Expected event numbers of the $2e$ and 2μ signals after implementing the kinematical cut. The event numbers are depicted as the functions of the compactification scale R with a fixed value $\delta/R = 10$ eV. The integrated luminosity is taken to be 30 fb^{-1}.

such that the effect of lepton flavor violation is minimized. It then turns out from (3.3) that all the constraints are satisfied for $\delta/R \geqslant 6.6$ eV. Finally, the SM Higgs mass is to be $m_h = 120$ GeV in evaluating the decay widths of heavy KK neutrinos ($N \to h + \nu$).

Now we are at the stage of investigating the tri-lepton signal of heavy neutrino productions at the LHC. Since the tau lepton is hardly detected compared to the others, we consider the signal event including only electrons and muons. There are four kinds of tri-lepton signals: eee, $ee\mu$, $e\mu\mu$, and $\mu\mu\mu$. In this work, we use two combined signals which are composed of $eee + ee\mu$ (the $2e$ signal) and $e\mu\mu + \mu\mu\mu$ (the 2μ signal). Fig. 2 shows the total cross sections for these signals from the 1st KK mode productions at the LHC. They are described as the functions of the compactification scale R with δ/R being 10 eV. It is found from the figure that the cross section for the 2μ signal is about one order of magnitude larger than the $2e$ signal.[3] We have also evaluated the cross sections of tri-lepton signals from heavier KK neutrinos and found that they are more than one order of magnitude smaller than the above and are out of reach of the LHC machine. A high luminosity collider with clean environment such as the International Linear Collider (ILC) would distinctly discover the signatures of KK mode resonances.

To clarify whether the tri-lepton signal is captured at the LHC, it is important to estimate SM backgrounds against the signal. The SM backgrounds which produce or mimic the tri-leptons final state have been studied [14,15], and for the present purpose a useful kinematical cut is discussed to reduce these SM processes [15]. According to that work, we adopt the following kinematical cuts;

- The existence of two like-sign charged leptons $\ell_1^{\pm}$, $\ell_2^{\pm}$, and an additional one with the opposite charge $\ell_3^{\mp}$.
- Both energies of the like-sign leptons are larger than 30 GeV.
- Both invariant masses from ℓ_1 and ℓ_3 and from ℓ_2 and ℓ_3 are larger than $m_Z + 10$ GeV or smaller than $m_Z - 10$ GeV.

The last condition is imposed to reduce large backgrounds from the leptonic decays of Z bosons in the SM processes. Fig. 3 shows the expected numbers of signal events after imposing the cuts stated above. The results are depicted by assuming the integrated luminosity 30 fb^{-1}. In order to estimate the efficiency for signal events due to the cuts, we have used the Monte Carlo simulation using the CalcHep code [16]. Since the event numbers of SM backgrounds after the cut are about 260 for the $2e$ signal and 110 for the 2μ signal [15], the 2μ events are expected to be observed if $1/R$ is less than a few hundred GeV.

The luminosity which is required to find the 2μ signal at the LHC is shown in Fig. 4 as a contour plot on the $(1/R, \delta/R)$ plane. The contour is obtained by computing the significance for the signal discovery,

$$S_{ig} \equiv \frac{S}{\sqrt{S+B}}, \tag{3.4}$$

[3] For the inverted hierarchy spectrum of light neutrinos, the $2e$ signal cross section becomes larger than the 2μ one.

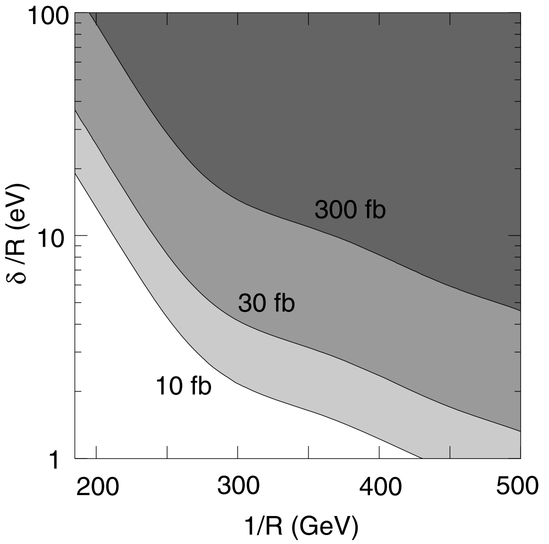

Fig. 4. Luminosity for the 3σ reach on the $(1/R, \delta/R)$ plane (10, 30, and 300 fb^{-1} contours).

where S and B are the numbers of the 2μ events and the corresponding SM backgrounds after the kinematical cut. Since both S and B are proportional to the luminosity, it is possible to estimate the luminosity, e.g. giving $S_{ig} = 3$ which is plotted in Fig. 4. The luminosity for signal confirmation (for $S_{ig} = 5$) are also obtained by scaling the above result. The luminosity of 10, 30, and 300 fb^{-1} are depicted in the figure. It is found that if $1/R$ is less than about 250 GeV, the signals will be observed at the early run of the LHC, while a larger luminosity is needed for a smaller size of extra dimension to find its signals.

4. Summary and discussion

We have discussed a seesaw scenario where right-handed neutrinos are around TeV scale, accessible in near future particle experiments. The seesaw-induced mass scale is of the order of eV, while the right-handed neutrinos have sizable gauge and Yukawa couplings to the SM sector. The scenario is a five-dimensional extension of the SM with right-handed neutrinos, where the ordinary SM particles locally live in four dimensions and the right-handed neutrinos exist in the bulk. The light neutrinos obtain tiny Majorana masses due to the small lepton number violation, and therefore the same-sign di-lepton processes cannot be observed. We have analyzed, as the most effective LHC signal, the lepton number preserving processes with tri-lepton final states, $pp \to \ell^\pm \ell^\pm \ell^\mp \nu(\bar{\nu})$. It is found that the scenario gives enough excessive tri-lepton events beyond the SM backgrounds in wide regions of parameter space, and the LHC would discover the signs of tiny neutrino mass generation and extra dimensions.

The possible experimental detections of neutrino mass generations have been discussed in other scenarios [4,9,17]. In the present analysis, only the 1st excited mode contributes to the signals. The observation of higher KK modes is expected to be within the reach of future collider experiments such as the ILC, which result makes the scenario substantially confirmed. Further analysis of such collider signatures, together with including bulk Dirac masses and curved gravitational backgrounds, are left for important future study.

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